

Top-quark and neutrino composite Higgs bosons

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In the context of top-quark condensation models, the top-quark alone is too light to saturate the correct value of the electroweak scale by its condensate. Within the seesaw scenario the neutrinos can have their Dirac masses large enough so that their condensates can provide significant contribution to the value of the electroweak scale. We address the question of a phenomenological feasibility of the top-quark and neutrino condensation conspiracy against the electroweak symmetry. Mandatory is to reproduce the masses of electroweak gauge bosons, the top-quark mass and the recently observed scalar mass of 125 GeV and to satisfy the upper limits on absolute value of active neutrino masses. To accomplish that we design a reasonably simplified effective model with two composite Higgs doublets. Additionally, we work with a general number N of right-handed neutrino flavor triplets participating on the seesaw mechanism. There are no experimental constraints limiting this number. The upper limit is set by the model itself. Provided that the condensation scale is of order 10^{17-18} GeV and the number of right-handed neutrinos is $\mathcal{O}(100 - 1000)$, the model predicts masses of additional Higgs bosons below 250 GeV and a suppression of the top-quark Yukawa coupling to the 125 GeV particle at the $\sim 60\%$ level of the Standard model value.

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I. INTRODUCTION

The Standard model reproduces the measured data surprisingly well. There is only little room for beyond-Standard model physics. The recently measured value for the Higgs boson mass $M_h \doteq 125$ GeV [1, 2] however leads to the running quartic coupling constant λ vanishing at the scale below the Planck scale [3–5]. There is conceivable interpretation: The Higgs boson reveals its compositeness at that scale. It corresponds to the historical experience that all of the observed scalar particles ultimately turned out to be fermion bound states.

The mechanism of the electroweak symmetry breaking manifests itself by the fact that both the top-quark mass m_t and the Higgs boson mass M_H are of the same order as the electroweak scale v which is the scale for mass of weak bosons, e.g., $M_W = g_2 v/2$. A simple contemplation leads to the suspicion that both the longitudinal components of the weak vector bosons and the Higgs boson are in fact bound states of top-quark. This is the idea of top-quark condensation models introduced in [6–8]. The key idea is that to break the electroweak symmetry it is enough to generate dynamically masses of known fermions, in particular of the top-quark. In this sense the top-quark mass generation is primary. The electroweak symmetry breaking comes out automatically. The top-quark mass then determines the magnitude of the electroweak scale and the Higgs boson mass.

Although the actual calculation gives a correct order of magnitude of masses, there are two essential failures of the simplest model when comparing it with experiment.

First, the top-quark is observed to be too light to saturate the electroweak scale v . Keeping the top-quark condensation scale below the Planck scale, the top-quark condensation alone can provide only at most 68 % of the W and Z boson masses. Second, the Higgs boson is predicted to be too heavy, in all available calculations $M_H > m_t$. This prediction was ruled out already before the actual measurement of 125 GeV particle at the LHC [1, 2].

If we want to maintain the attractive idea that the top-quark condensation is the source of the electroweak symmetry breaking, we need to improve the simplest model.

First, we need to suppose yet another source potent to saturate the value of the electroweak scale. The remaining observed quarks and charged leptons bring no improvement in this respect as they are way too light and contribute by truly negligible amount to the electroweak scale. At this point one could be easily seduced to invoke some new yet unobserved fermions with high enough mass, like, for instance, fourth-generation fermions or techniquarks. This however may not be necessary. Among the known fermions, there is an additional source of the electroweak symmetry breaking naturally present in the form of the neutrino Dirac mass m_D provided that the seesaw mechanism is in work. If the neutrino Dirac mass is of the order of the electroweak scale, $m_D \sim v$, then the neutrino condensate is strong enough to complement the electroweak scale [9–11].

Second, once we identify two main fermion sources of the electroweak symmetry breaking we resort to the reasonably simplified description using correspondingly two composite Higgs doublets. The presence of more than single Higgs doublet is instrumental in achieving lighter neutral scalar as the candidate for the 125 GeV particle due to the appropriate mixing among components of the doublets.

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The idea of top-quark and neutrino condensation was addressed already in the past. First, Martin [9] investigated the model in which the idea was implemented in the simplest possible way. He invoked a factorization assumption on four-fermion interactions which resulted in the low-energy description with only single Higgs doublet. He reached the correct value of top-quark mass, but from now-days perspective, the model suffers from exhibiting too heavy Higgs boson particle, in the same way as the original top-quark-alone condensation models [7, 8]. Ten years later, the issue was addressed again by Antusch et al. confirming its usefulness for obtaining the correct value of top-quark and suggesting that the two-Higgs doublet low-energy description is worthwhile to study in more detail. Another ten years later we are addressing the idea once again and confront it with the new experimental evidence of the 125 GeV boson excitation.

We believe that analysing the top-quark and neutrino condensation is relevant for the class of models where all types of known fermion matter condense [12–24]. The models share the underlying principle of primary role of known fermion masses in the electroweak symmetry breaking. Those models which provide the seesaw mechanism for neutrino masses, like for instance [10, 25], have a chance to escape from being ruled out by the measurement of the top-quark mass, just because of the presence of the second potentially big Dirac mass m_D . Currently, however, they have to be confronted with the measurement of 125 GeV particle. We do not refer to any of these underlying models in our analysis, we rather emulate their dynamics by particular four-fermion interaction. Effectively, instead of using a multitude of composite Higgs doublets, one for each Dirac mass, we work with only two Higgs doublets, one for each significant source of the electroweak symmetry breaking.

The number of right-handed neutrino types participating on the seesaw mechanism is not constrained by any upper limit [26, 27]. As claimed there, higher number of the order $\mathcal{O}(100)$ is even well motivated within some string constructions. Large number of right-handed neutrinos $\mathcal{O}(100)$ has also an improving effect on the standard thermal leptogenesis [28]. Being of order $\mathcal{O}(10 - 100)$ it can serve as the reason for large lepton mixing angles [29]. Therefore we study the dependence of our results on the number of right-handed neutrinos.

The precise form of the neutrino mass spectrum is not known. For our analysis of electroweak symmetry breaking it does not play an essential role. Therefore we simulate the unknown mass spectrum of neutrinos by the most simple choice for the neutrino mass matrix, which is characterized by a common Dirac mass m_D , by a common right-handed Majorana mass M_R and by the number N of right-handed neutrino flavor triplets. By this simplification we can control the order of magnitude of active neutrino masses but do not reproduce the neutrino physics exactly, what would otherwise require to specify in detail the underlying dynamics.

In this work we explore the possibility to saturate the electroweak symmetry breaking by combined sources from both top-quark and neutrino Dirac condensates. We concentrate on the analysis of the low-energy mass spectrum, mainly to the possibility to achieve one of the composite scalar as light as 125 GeV while reproducing correct values for the masses of electroweak gauge bosons and top-quark and satisfying the upper limits on absolute value of active neutrino masses $m_\nu < 0.2$ eV. We study coupling properties of the lighter Higgs scalar to the top-quark and gauge bosons. Next we calculate the mass spectrum of additional Higgs bosons. We study the sensitivity of the results on the number N of right-handed neutrino flavor triplets.

In the Sec. II we introduce the lagrangian for our analysis and identify relevant symmetries. In the Sec. III we formulate the effective description using two Higgs doublets and write the low-energy lagrangian. In the Sec. IV we study the mass spectrum of the model as a function of parameters of the low-energy lagrangian. In the Sec. V we formulate the renormalization group equations which govern the evolution of the low-energy parameters from the condensation scale down to the electroweak scale. In the Sec. VI we present results. In the Sec. VII we discuss the results and confront them with experimental data. Finally, in the Sec. VII we conclude.

II. UNDERLYING FOUR-FERMION INTERACTION

A. Underlying lagrangian

For the purpose of our analysis we let only top-quark and neutrinos to condense and to contribute to the electroweak scale. Therefore we define our simplified model by the four-fermion interaction

$$\mathcal{L}_{4f} = -G_t(\bar{t}_R q_L)(\bar{q}_L t_R) - G_\nu \left(\sum_s \bar{\nu}_{Rs} \ell_L \right) \left(\sum_{s'} \bar{\ell}_L \nu_{Rs'} \right) \quad (1)$$

which is designed to provide us just by the condensation of the desired form discussed in the introduction. It reflects the main assumption that only the top-quark and neutrinos play the appreciable role in what we address in our work - the electroweak symmetry breaking.

Only the third generation of quarks participates in the interaction. The fields t_R and q_L are color triplets. On the other hand, because we suppose that all neutrino Dirac masses are of the order of electroweak scale then within the simplified model we are letting all three generations of leptons to participate in the interaction. Therefore the fields ν_R and ℓ_L are flavor triplets. The left fields are weak isospin doublets. All three types of indices are suppressed. The explicitly written index $s = 1, \dots, N$ labels N right-handed neutrino flavor triplets. By this simplified dynamics we are going to generate masses of only top-quark and neutrinos.

If the underlying dynamics is such that the four-fermion interactions follow from an exchange of neutral and colorless gauge bosons, then there are only these two effective terms relevant for the top-quark and neutrino condensation. No mixing terms like $\propto (\bar{t}_R q_L)(\bar{\ell}_L \nu_R)$ appear. There could appear also various four-fermion interactions of other leptons and quarks, but we neglect them here as they play rather negligible role in the electroweak symmetry breaking.

For the sake of simplicity we use here a factorization assumption on the neutrino coupling constants, so that there is single neutrino coupling constant G_ν and single neutrino-Higgs boson doublet, in the same spirit as in [9]. After the condensation, this provides a degenerate Dirac mass for neutrinos.

We introduce the right-handed neutrino Majorana mass term. We take it degenerate and diagonal for the same sake of simplicity

$$\mathcal{L}_{M_R} = -\frac{1}{2} M_R \overline{\nu_{Rs}^c} \nu_{Rs} + \text{h.c.} \quad (2)$$

We dare to accept these simplifications as our aim is not to reproduce the neutrino phenomenology exactly.

For our purpose we assume the lagrangian

$$\mathcal{L} = \mathcal{L}_{\text{usual}} + \mathcal{L}_{\text{model}}, \quad (3)$$

$$\mathcal{L}_{\text{model}} = \mathcal{L}_{4f} + \mathcal{L}_{M_R}, \quad (4)$$

where $\mathcal{L}_{\text{usual}}$ contains kinetic terms of all known fermions, their Standard model gauge interactions and pure gauge boson terms.

B. Symmetries

The lagrangian $\mathcal{L}_{\text{model}}$ has well separated quark and lepton sectors. On the classical level, it is invariant under the global symmetry¹

$$G_{\text{model}} = [\text{SU}(2) \times \text{U}(1)^2]_q \times [\text{SU}(2) \times \text{U}(1)]_\ell. \quad (5)$$

One subgroup of G_{model} is the electroweak $\text{SU}(2)_L \times \text{U}(1)_Y$ gauge symmetry group. The electroweak interactions explicitly break the symmetry G_{model} , so the symmetry of the full lagrangian \mathcal{L} is

$$G = \text{SU}(2)_L \times \text{U}(1)_Y \times \text{U}(1)_B \times \text{U}(1)_X, \quad (6)$$

among which $\text{SU}(2)_L \times \text{U}(1)_Y$ are the weak isospin and weak hypercharge gauge symmetries, $\text{U}(1)_B$ is the baryon number and $\text{U}(1)_X$ is the axial symmetry

$$\text{U}(1)_X : (q_L, t_R, \ell_L, \nu_R) = (-1, 0, 1, 0). \quad (7)$$

¹ The rest of standard fermions and their corresponding symmetries are of course present in the model in order to provide proper anomaly cancelation, but we hold them back here as they do not participate in the symmetry breaking in our simplified analysis. Due to the factorization assumption the three generations of leptons exhibit a single common symmetry group.

We are making this choice of X charges in order to have $X(\bar{t}_R q_L) = -1$ and $X(\bar{\nu}_R \ell_L) = +1$. It is this symmetry which prevents the top-quark and neutrino sectors from mixing.

On the quantum level, the group $\text{U}(1)_X$ is axially anomalous. The axial anomaly is provided by the electroweak and QCD dynamics. Additionally, it can be provided by some new not specified dynamics underlying the four-fermion interaction (1) like, e.g., the gauge flavor dynamics [24]. Our ignorance about the anomaly stems from not specifying the complete structure of the model in order to be flexible in our analysis. In the following, we will simply parameterize the strength of the anomaly by the scale $\mu_{t\nu}$.

The dynamically generated Dirac masses for top-quark and neutrinos break spontaneously the G_{model} symmetry down to

$$G_{\text{model}} \xrightarrow{m_t, m_\nu} [\text{U}(1)^2]_q \times [\text{U}(1)]_\ell. \quad (8)$$

It would give rise to 6 massless Nambu–Goldstone bosons.

There are however effects of both the electroweak dynamics and of the axial anomaly which eliminate the 6 massless states completely. The electroweak interactions change the spontaneous symmetry breaking pattern to

$$G \xrightarrow{m_t, m_\nu} \text{U}(1)_{\text{em}} \times \text{U}(1)_B. \quad (9)$$

Therefore three of the Nambu–Goldstone states are eaten by the electroweak gauge bosons. The other two form a single charged pseudo-Nambu–Goldstone particles whose mass results from the explicit breaking by the electroweak dynamics (6) and it is therefore proportional to the electroweak gauge coupling constants. The remaining single state stays massless if we neglect the effect of the $\text{U}(1)_X$ axial anomaly, otherwise it is the pseudo-Nambu–Goldstone boson with the mass proportional to the scale $\mu_{t\nu}$.

III. TWO HIGGS DOUBLET DESCRIPTION

Effectively, the top-quark and neutrino condensation can be described by the condensation of two composite Higgs doublets

$$H_t \sim (\bar{t}_R q_L), \quad (10)$$

$$H_\nu \sim \left(\sum_s \bar{\nu}_{Rs} \ell_L \right). \quad (11)$$

Using them we can rewrite the four-fermion interaction as

$$\begin{aligned} \mathcal{L}_{4f} = & -y_{t0}(\bar{q}_L t_R)H_t - y_{\nu 0} \left(\sum_s \bar{\ell}_L \nu_{Rs} \right) H_\nu + \text{h.c.} \\ & + \mu_{t0}^2 H_t^\dagger H_t + \mu_{\nu 0}^2 H_\nu^\dagger H_\nu, \end{aligned} \quad (12)$$

which is completely equivalent to (1) as far as H_t and H_ν are non-propagating auxiliary fields.

Below the condensation scale Λ , we substitute the original lagrangian for the effective lagrangian formed by operators allowed by the symmetries. Among the operators, there are kinetic terms for the composite Higgs doublets H_t and H_ν . If we take only the renormalizable operators we are effectively obtaining two-Higgs-doublet model.

$$\mathcal{L}_{\text{eff}} = |DH_t|^2 + |DH_\nu|^2 - \mathcal{V}(H_t, H_\nu) - y_t(\bar{q}_L t_R)H_t - y_\nu(\sum_s \bar{\ell}_L \nu_{Rs})H_\nu + \text{h.c.} \quad (13)$$

The potential for the two Higgs doublets invariant with respect to the G symmetry is

$$\begin{aligned} \mathcal{V} &= \mathcal{V}_0 + \mathcal{V}_{\text{EW}} + \mathcal{V}_{\text{soft}} \\ \mathcal{V}_0 &= -\mu_t^2 H_t^\dagger H_t - \mu_\nu^2 H_\nu^\dagger H_\nu \\ &\quad + \frac{1}{2}\lambda_t (H_t^\dagger H_t)^2 + \frac{1}{2}\lambda_\nu (H_\nu^\dagger H_\nu)^2 \\ \mathcal{V}_{\text{EW}} &= \lambda_{t\nu} (H_t^\dagger H_t)(H_\nu^\dagger H_\nu) + \lambda'_{t\nu} (H_t^\dagger H_\nu)(H_\nu^\dagger H_t). \end{aligned} \quad (14)$$

We sort the terms in the potential for Higgs bosons according to their primary origin. Those terms denoted by \mathcal{V}_0 are generated due to the four-fermion interaction irrespectively of the presence of the electroweak dynamics which provides only corrections to their magnitude. The terms denoted by \mathcal{V}_{EW} on the other hand make a bridge between the top-quark and neutrino sectors generated only because of the presence of the electroweak dynamics. They vanish in the limit of vanishing electroweak coupling constants.

In order to take into account the axial anomaly of $U(1)_X$ we introduce additional term which mixes the two Higgs doublets and breaks explicitly the $U(1)_X$ symmetry

$$\mathcal{V}_{\text{soft}} = -\mu_{t\nu}^2 H_t^\dagger H_\nu. \quad (15)$$

This term can not be generated at any loop order neither by the four-fermion interaction nor by the electroweak dynamics. In this work we use $\mu_{t\nu}$ as a free parameter.

Apart from $\mu_{t\nu}$ all the parameters of the lagrangian \mathcal{L}_{eff} , i.e., y 's, μ 's, and λ 's, run with the renormalization scale μ according to the renormalization group equations towards the condensation scale $\mu = \Lambda$. At the condensation scale they are linked to the values of the underlying lagrangian \mathcal{L} via the field renormalization factors. Because of its non-perturbative origin, the mixing parameter $\mu_{t\nu}$ is not the subject of the renormalization group equations and as such it acts as a free parameter.

The quartic stability of the potential is given by the conditions

$$\lambda_t, \lambda_\nu > 0, \quad (16)$$

$$\sqrt{\lambda_t \lambda_\nu} > -\lambda_{t\nu} - \lambda'_{t\nu}. \quad (17)$$

The parameter setting in the range

$$\lambda'_{t\nu} < 0 \text{ and } \mu_{t\nu}^2 > 0 \quad (18)$$

leads to the minimum of the potential which conserves electric charge. This is completely analogous to the work [30] in the context of top-quark and bottom-quark two-Higgs-doublet model.

IV. MASSES

A. Electroweak scale and fermion masses

The top-quark and the neutrino contributions to the electroweak symmetry breaking are given by the condensates $v_t \propto \langle \bar{t}t \rangle$, and $v_\nu \propto \langle \bar{\nu}\nu \rangle$.

From the lagrangians (2) and (13) the neutrino mass matrix \mathbf{M}_ν turns out to be

$$\mathbf{M}_\nu = \begin{pmatrix} 0 & \mathbf{m}_D & \cdots & \mathbf{m}_D \\ \mathbf{m}_D & \mathbf{m}_R & & \\ \vdots & & \ddots & \\ \mathbf{m}_D & & & \mathbf{m}_R \end{pmatrix}, \quad (19)$$

where $\mathbf{m}_D \equiv \frac{1}{\sqrt{2}} y_\nu v_\nu \mathbb{1}$ and $\mathbf{m}_R \equiv M_R \mathbb{1}$, where $\mathbb{1}$ is the 3×3 unit matrix in the flavor space. Each left-handed neutrino therefore mixes with N right-handed neutrinos, but do not mix among each other.

We assume that both v_t and v_ν together saturate the electroweak scale v ,

$$v^2 = v_t^2 + v_\nu^2. \quad (20)$$

This is the main idea of this work.

Next we define β -angle by

$$\tan \beta \equiv \frac{v_t}{v_\nu}. \quad (21)$$

We choose the convention that $\beta \in (0, \frac{\pi}{2})$.

The mass of top-quark is given by equation

$$m_t = y_t(\mu = m_t) v_t / \sqrt{2} \quad (22)$$

and the light neutrino mass is given as a smaller eigenvalue of the neutrino mass matrix (19) by the seesaw formula

$$m_\nu = \frac{N y_\nu^2(\mu = m_\nu) v_\nu^2}{2M_R}, \quad (23)$$

where $y_{t,\nu}(\mu)$ are running Yukawa coupling constants.

Within the two-composite-Higgs-doublet model the values of Yukawa couplings around the electroweak scale are calculable using their renormalization group evolution down from the condensation scale Λ . Because Λ is very large the couplings are only weakly sensitive to their initial values $y_{t,\nu}(\Lambda)$ as they have enough time to approach an infrared fixed point.

The equation (22) for the top-quark mass, $m_t \doteq 172 \text{ GeV}$, then fixes v_t . From the relation (20) we determine v_ν – a portion of the electroweak scale left for neutrinos. Finally from (23) we determine the right-handed neutrino Majorana mass M_R based on the assumption that $m_\nu \lesssim 0.2 \text{ eV}$. Having $|\mathbf{m}_D| \sim v$ implies roughly $M_R \gtrsim 10^{14} \text{ GeV}$. Here the construction closes by an important restriction on the right-handed neutrino Majorana mass and the condensation scale

$$\Lambda_{\text{Planck}} > \Lambda > M_R. \quad (24)$$

If the Majorana mass M_R were too big, then the correspondingly heavy right-handed neutrinos would decouple from the dynamics [9] before they would manage to condense with the left-handed neutrinos. Finally, we find the top-quark and neutrino condensation meaningful only if the condensation scale Λ is below the Planck scale.

B. Higgs boson masses

After the electroweak symmetry is broken by the condensate $v^2 = v_t^2 + v_\nu^2$, we get Higgs boson mass eigenstates from \mathcal{L}_{eff} . The masses for neutral scalars H and h , for the neutral pseudo-scalar A and for the charged scalar H^\pm are given by

$$M_{H,h}^2 = \frac{1}{2} f_\pm(t = \ln M_{H,h}), \quad (25a)$$

$$M_A^2 = \frac{2\mu_{t\nu}^2}{\sin 2\beta}, \quad (25b)$$

$$M_{H^\pm}^2 = \frac{2\mu_{t\nu}^2}{\sin 2\beta} - \frac{1}{2} \lambda'_{t\nu}(t = \ln M_{H^\pm}) v^2 \quad (25c)$$

where

$$\begin{aligned} f_\pm(t) &= v_t^2 \lambda_t(t) + v_\nu^2 \lambda_\nu(t) + \frac{2\mu_{t\nu}^2}{\sin 2\beta} \pm \sqrt{A(t)}, \\ A(t) &= (v_t^2 \lambda_t(t) - v_\nu^2 \lambda_\nu(t))^2 + 4v_t^2 v_\nu^2 (\lambda_{t\nu}(t) + \lambda'_{t\nu}(t))^2 \\ &\quad + \frac{4\mu_{t\nu}^4}{\sin^2 2\beta} - 2\mu_{t\nu}^2 [v_t^2 \lambda_t(t) \tan \beta + v_\nu^2 \lambda_\nu(t) \cot \beta \\ &\quad - v_t v_\nu (\lambda_t(t) + \lambda_\nu(t) + 4\lambda_{t\nu}(t) + 4\lambda'_{t\nu}(t))]. \end{aligned}$$

We have introduced the logarithmic renormalization scale $t \equiv \ln \mu$.

The mixing of the two neutral scalars is given by the angle α

$$H = \sqrt{2} [\text{Re } H_t^0 \sin \alpha + \text{Re } H_\nu^0 \cos \alpha], \quad (26)$$

$$h = \sqrt{2} [\text{Re } H_t^0 \cos \alpha - \text{Re } H_\nu^0 \sin \alpha], \quad (27)$$

$$\tan 2\alpha = \frac{(\lambda_{t\nu} + \lambda'_{t\nu}) v_t v_\nu - \mu_{t\nu}^2}{-\frac{1}{2}(v_t^2 \lambda_t - v_\nu^2 \lambda_\nu) - \mu_{t\nu}^2 \cot 2\beta}. \quad (28)$$

We choose the convention that $\alpha \in (-\frac{\pi}{2}, 0)$. In this case the lighter Higgs scalar is always h .

In the following we will show that our model leads to

$$\lambda_{t\nu}, \lambda'_{t\nu} < 0 \quad (29)$$

and

$$\lambda_t, \lambda_\nu \gg |\lambda_{t\nu}|, |\lambda'_{t\nu}|. \quad (30)$$

For illustration, see Fig. 1. In order to conserve electric charge we will investigate only the values

$$\mu_{t\nu}^2 \geq 0. \quad (31)$$

Let us investigate two important limits:

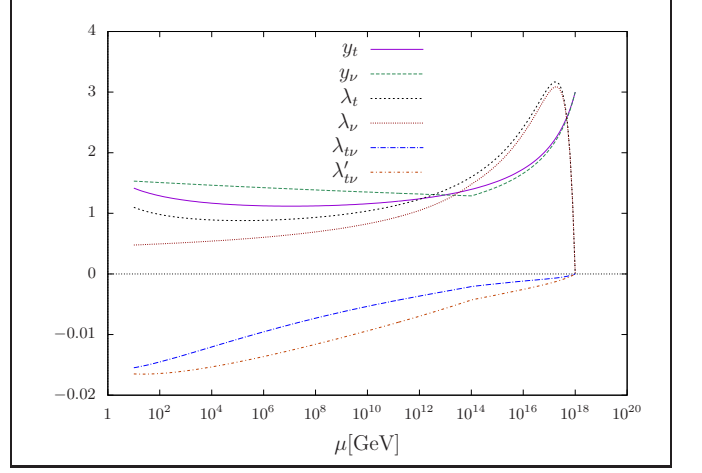


FIG. 1: The renormalization group evolution of y_t , y_ν , λ_t , λ_ν , $\lambda_{t\nu}$ and $\lambda'_{t\nu}$ for the parameter setting 47. The role of $M_R = 10^{14}$ GeV is visible as a threshold in the evolution.

1) First, by setting $\lambda_{t\nu}, \lambda'_{t\nu} \rightarrow 0$, we switch the electroweak interactions off, and by setting $\mu_{t\nu}^2 = 0$, we preserve the $U(1)_X$ symmetry. The top-quark and neutrino condensation causes the symmetry breaking (9). In this limit the spectrum of bosons changes to

$$M_A^2 = M_{H^\pm}^2 = 0, \quad (32)$$

$$M_H = \text{Max}\{v_t^2 \lambda_t, v_\nu^2 \lambda_\nu\}, \quad (33)$$

$$M_h = \text{Min}\{v_t^2 \lambda_t, v_\nu^2 \lambda_\nu\}, \quad (34)$$

where we identify three uneaten Nambu–Goldstone bosons. The two Higgs scalars do not mix as $\alpha = -\frac{\pi}{2}$. The identity of the lighter Higgs scalar is given by

$$\text{for } v_t^2 \lambda_t > v_\nu^2 \lambda_\nu : \quad h = \sqrt{2} \text{Re } H_\nu^0, \quad (35)$$

$$\text{for } v_t^2 \lambda_t < v_\nu^2 \lambda_\nu : \quad h = \sqrt{2} \text{Re } H_t^0. \quad (36)$$

2) Second, we set $\lambda_{t\nu}, \lambda'_{t\nu} \rightarrow 0$ again, but we let $\mu_{t\nu} \gg v$. In this limit, the spectrum of bosons changes to

$$M_H = M_A^2 = M_{H^\pm}^2 = \frac{2\mu_{t\nu}^2}{\sin 2\beta}, \quad (37)$$

$$\begin{aligned} M_h &= \frac{1}{8} v^2 [4\lambda_t (\sin^2 \beta + \sin^4 \beta) + 4\lambda_\nu (\cos^2 \beta + \cos^4 \beta) \\ &\quad - (\lambda_t + \lambda_\nu) \sin^2 2\beta]. \end{aligned} \quad (38)$$

In this limit four degrees of freedom H , A and H^\pm get degenerate masses proportional to $\mu_{t\nu}$ and decouple from the low-energy physics. One degree of freedom h stays light. It is the mixture of top-quark and neutrino neutral composite scalars which is characterized by $\tan 2\alpha = \tan 2\beta$, hence the mixing angle $\alpha = \beta - \frac{\pi}{2}$.

To obtain values of masses we need to evolve the running parameters from their initial values at the condensation scale down to the electroweak scale. This we will do in the next section.

C. Interactions of mass eigenstates

We list here several interactions of the lighter Higgs scalar h and of the charged Higgs boson H^\pm , which are particularly important for today phenomenological analysis:

$$\mathcal{L}_{\text{Yukawa}}^h = -\frac{m_t}{v} C_t h \bar{t} t, \quad (39a)$$

$$\mathcal{L}_{\text{gauge}}^h = -C_V h \left(\frac{2M_W^2}{v} W^+ W^- + \frac{M_Z^2}{v} Z^2 \right) \quad (39b)$$

$$\mathcal{L}_{H^\pm}^h = -C_{H^\pm} v h H^+ H^-, \quad (39c)$$

where the scaling coupling factors are

$$C_t = \frac{\cos \alpha}{\sin \beta}, \quad (40a)$$

$$C_V = \sin(\beta - \alpha), \quad (40b)$$

$$C_{H^\pm} = \sin \beta \cos \alpha (\sin^2 \beta \lambda_{t\nu} + \cos^2 \beta (\lambda_t - \lambda'_{t\nu})) - \cos \beta \sin \alpha (\cos^2 \beta \lambda_{t\nu} + \sin^2 \beta (\lambda_\nu - \lambda'_{t\nu})). \quad (40c)$$

They measure the departure from the Standard model, which is characterized by $C_t = 1$, $C_V = 1$ and $C_{H^\pm} = 0$.

V. RENORMALIZATION GROUP EQUATIONS

The parameters of the low-energy lagrangian \mathcal{L}_{eff} run with the renormalization scale μ according to equations of renormalization group. The exhausting analysis of renormalization group equations for two-Higgs-doublet models is to be found in [31]. The compositeness of the Higgs doublets is expressed by the fact that the lagrangian (13) is equivalent to the Higgs-less lagrangian (1) or (12) at the condensation scale Λ . From sewing the two lagrangian together at Λ , a set of boundary conditions for $\mu \rightarrow \Lambda$ follows

$$\begin{aligned} y_t \rightarrow \infty, \quad y_\nu \rightarrow \infty, \quad y_t/y_\nu \rightarrow y_{t0}/y_{\nu 0}, \\ \lambda_t/y_t^4 \rightarrow 0, \quad \lambda_\nu/y_\nu^4 \rightarrow 0, \\ \lambda_{t\nu}/y_t^2 y_\nu^2 \rightarrow 0, \quad \lambda'_{t\nu}/y_t^2 y_\nu^2 \rightarrow 0, \\ \mu_t^2/y_t^2 \rightarrow \mu_{t0}^2/y_{t0}^2, \quad \mu_\nu^2/y_\nu^2 \rightarrow \mu_{\nu 0}^2/y_{\nu 0}^2, \end{aligned} \quad (41)$$

In practice, for actual numerical calculation, we will use the boundary conditions

$$\begin{aligned} y_t(\ln \Lambda) &= Y_t, \quad y_\nu(\ln \Lambda) = Y_\nu, \\ \lambda_t(\ln \Lambda) &= 0, \quad \lambda_\nu(\ln \Lambda) = 0, \\ \lambda_{t\nu}(\ln \Lambda) &= 0, \quad \lambda'_{t\nu}(\ln \Lambda) = 0, \end{aligned} \quad (42)$$

where Y_t and Y_ν are finite numbers on which the low-energy result depends only very weakly.

Further, we will restrict our analysis only to one loop order.

The presence of the second Higgs doublet affects the t evolution of the gauge coupling constants governed by the one-loop renormalization group equations and the boundary conditions given by the experimental values at $\mu = M_Z$:

$$16\pi^2 \frac{d}{dt} g_1 = 7g_1^3, \quad g_1^2(\ln M_Z) \doteq 0.127, \quad (43a)$$

$$16\pi^2 \frac{d}{dt} g_2 = -3g_2^3, \quad g_2^2(\ln M_Z) \doteq 0.425, \quad (43b)$$

$$16\pi^2 \frac{d}{dt} g_3 = -7g_3^3, \quad g_3^2(\ln M_Z) \doteq 1.440. \quad (43c)$$

The renormalization group equation for Yukawa coupling constants are

$$16\pi^2 \frac{d}{dt} y_t = y_t \left[\frac{9}{2} y_t^2 - \frac{17}{12} g_1^2 - \frac{9}{4} g_2^2 - 8g_3^2 \right], \quad (44)$$

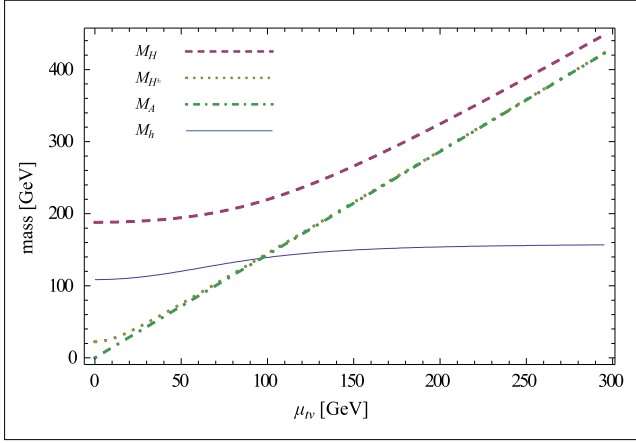
$$16\pi^2 \frac{d}{dt} y_\nu = y_\nu \left[3(N + \frac{1}{2})\theta(t - M_R)y_t^2 - \frac{3}{4}g_1^2 - \frac{9}{4}g_2^2 \right].$$

The θ -function stands for threshold under which the heavy right-handed neutrinos decouple from the system. The renormalization group equations for the quartic coupling constants are

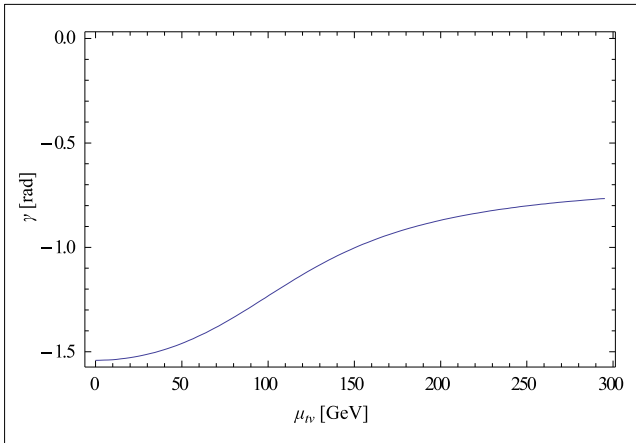
$$\begin{aligned} 16\pi^2 \frac{d}{dt} \lambda_t &= 12\lambda_t^2 + 4\lambda_{t\nu}^2 + 4\lambda_{t\nu}^2 \lambda'_{t\nu} + 2\lambda_{t\nu}'^2 + \lambda_t [12y_t^2 - 3g_1^2 - 9g_2^2] - 12y_t^4 + \frac{3}{4}(g_1^4 - 2g_1^2 g_2^2 + 3g_2^4), \\ 16\pi^2 \frac{d}{dt} \lambda_\nu &= 12\lambda_\nu^2 + 4\lambda_{t\nu}^2 + 4\lambda_{t\nu}^2 \lambda'_{t\nu} + 2\lambda_{t\nu}'^2 + \lambda_\nu [12N\theta(t - M_R)y_\nu^2 - 3g_1^2 - 9g_2^2] - 12Ny_\nu^4 + \frac{3}{4}(g_1^4 - 2g_1^2 g_2^2 + 3g_2^4), \\ 16\pi^2 \frac{d}{dt} \lambda_{t\nu} &= 2(\lambda_\nu + \lambda_t)(3\lambda_{t\nu} + \lambda'_{t\nu}) + 4\lambda_{t\nu}^2 + 2\lambda_{t\nu}'^2 + \lambda_{t\nu} [6y_t^2 + 6N\theta(t - M_R)y_\nu^2 - 3g_1^2 - 9g_2^2] + \frac{3}{4}(g_1^4 - 2g_1^2 g_2^2 + 3g_2^4), \\ 16\pi^2 \frac{d}{dt} \lambda'_{t\nu} &= \lambda'_{t\nu} [\lambda_\nu + \lambda_t + 8\lambda_{t\nu}^2 + 4\lambda_{t\nu}'^2 + 6y_t^2 + 6N\theta(t - M_R)y_\nu^2 - 3g_1^2 - 9g_2^2] + 3g_1^2 g_2^2. \end{aligned} \quad (45)$$

According to Luty [30] the one loop renormalization evolution of the dimensionless parameters does not depend

on the presence of the mixing parameter $\mu_{t\nu}$.



a)



b)

FIG. 2: a) Higgs masses $M_{H,h}$, M_A and M_{H^\pm} as functions of $\mu_{t\nu}$ for $\Lambda = 10^{18}$ GeV, $N = 1$ and fixed neutrino mass to be $m_\nu = 0.2$ eV (i.e. $M_R \sim 10^{14}$ GeV). b) Mixing of two scalar Higgs bosons H and h as a function of $\mu_{t\nu}$. At the condensation scale, the Yukawa couplings $y_{t,\nu}$ start at the value $Y_{t,\nu} = 3$, and the quartic couplings λ start at zero value.

VI. RESULTS

Before we present our results we briefly describe the strategy of the analysis. Our input parameters are Λ , M_R and $\mu_{t\nu}$, out of which M_R is fixed by reproducing the neutrino mass $m_\nu = 0.2$ eV. Strictly speaking, there are additional two parameters Y_t and Y_ν which however have only very mild effect on the result and we take them quite arbitrarily to be

$$Y_t = Y_\nu = 3. \quad (46)$$

On top of that we have freedom to choose the number N of right-handed neutrino triplets.

The steps of calculation follow: 1) We solve analytically the equations (43) for gauge constants. 2) We nu-

merically evolve the Yukawa and quartic coupling constants according to the equations (44) and (45) with the boundary conditions (42) and (46). 3) From the experimental value of the top quark mass we determine v_t using the equation (22). 4) We calculate v_ν and β from equations (20) and (21), respectively. 5) The equation (23) gives us a value for the neutrino mass m_ν . Changing the value of M_R we repeat the calculation from the point 2) and iterate the value of neutrino mass to get $m_\nu = 0.2$ eV. 6) Using the equations (25) we calculate the Higgs boson mass spectrum.

A. Renormalization group evolution

The renormalization group evolution of the dimensionless parameters are plotted in Fig. 1 for

$$N = 1, \Lambda = 10^{18} \text{ GeV}, M_R = 10^{14} \text{ GeV}. \quad (47)$$

As we mentioned before, at one-loop order the result does not depend on the parameter $\mu_{t\nu}$. It gives us typical result which does not change qualitatively much with changing Λ and M_R . We can see that at the electroweak scale Λ_{EW}

$$\lambda_{t\nu}, \lambda'_{t\nu} < 0, \quad (48)$$

$$100 \times (|\lambda_{t\nu}|, |\lambda'_{t\nu}|) \sim (\lambda_t, \lambda_\nu), \quad (49)$$

so the stability conditions (29) and (30) are fulfilled. On top of that, taking $\mu_{t\nu}^2 > 0$ we assure that the condensate will be electrically neutral.

B. Mass spectrum of Higgs bosons

The typical result (49) allows us to neglect $\lambda_{t\nu}$, $\lambda'_{t\nu}$ in favor of λ_t , λ_ν . That is why the limits analyzed at the end of Sec. IV B are useful for us. We can roughly estimate the mass of the lighter Higgs scalar h and the mixing angle α to lie in the interval

$$\left. \begin{aligned} M_h &\simeq (113, 160) \text{ GeV} \\ \alpha &\simeq (-\frac{\pi}{2}, -0.7) \end{aligned} \right\} \text{ for } \mu_{t\nu} = (0, \infty) \text{ GeV} \quad (50)$$

calculated from (34) and (38) with input parameters (47) and using estimated values from Fig. 1, (20) and (22)

$$v_t \simeq 187 \text{ GeV}, v_\nu \simeq 160 \text{ GeV}, \quad (51)$$

$$\lambda_t \simeq 1.0, \lambda_\nu \simeq 0.5. \quad (52)$$

It represents a promising improvement with respect to the previous results of the single Higgs doublet top-quark condensation models (for review see [32]).

Now, let us investigate the solutions of the Higgs boson mass spectrum without approximations. In Fig. 2a) we plot the dependence of Higgs boson masses on the mixing parameter $\mu_{t\nu}$. We use $\Lambda = 10^{18}$ GeV while the right-handed neutrino Majorana mass we fix from demanding $m_\nu = 0.2$ eV to be roughly $M_R \sim 10^{14}$ GeV.

For lower values of $\mu_{t\nu}$ the bosons A and H^\pm are very light, the mass M_A even vanishes for vanishing $\mu_{t\nu}$ reflecting the spontaneous symmetry breaking of the $U(1)_X$ symmetry (7). In Fig. 2b) where we plot the dependence of the h - H mixing angle α on $\mu_{t\nu}$, it can be seen that the lighter scalar h is composed mainly by neutrinos, $\alpha \sim -\frac{\pi}{2}$, and therefore its coupling to top-quark is suppressed, see (39a).

Increasing $\mu_{t\nu}$ translates into the lifting of masses of the Higgs bosons, H , A and H^\pm . They soon become growing linearly and nearly degenerate. On the other hand, the mass of the lighter scalar h is only mildly sensitive to the increase of $\mu_{t\nu}$ and quite soon saturates just below 160 GeV. On top of that it is acquiring gradually larger admixture from a top-quark composite state, reaching the value over $\alpha \sim -0.8$. M_A minimizes the spectrum of H , A and H^\pm for all positive values of $\mu_{t\nu}$ in our model.

Setting $\Lambda = 10^{18}$ GeV, $m_\nu = 0.2$ eV and $N = 1$ we can reach the lighter Higgs boson mass of the desired value

$$M_h = 125 \text{ GeV for } \mu_{t\nu} \doteq 62 \text{ GeV.} \quad (53)$$

The value $\mu_{t\nu} \doteq 62$ GeV translates into the Higgs boson mass spectrum

$$M_H \doteq 198 \text{ GeV, } M_A \doteq 88 \text{ GeV, } M_{H^\pm} \doteq 91 \text{ GeV.} \quad (54)$$

These values can be altered by changing the model parameters, the condensation scale Λ and the number of right-handed neutrino triplets N .

By an order of magnitude decrease of Λ for fixed number N and $M_h = 125$ GeV we decrease the parameter $\mu_{t\nu}$ and also change the value of $\tan\beta$ according to Tab. I. The values for Λ above the Planck scale are shown only for curiosity, otherwise we avoid them further in our analysis.

By increasing the number of right-handed neutrino triplets N for a given Λ we can increase the value of $\mu_{t\nu}$. On the other hand the value of $\tan\beta$ is completely insensitive to the change of N .

Surprisingly, the number N has an upper limit given by either of two conditions: the non-decoupling condition $\Lambda > M_R$ (24), or the Higgs potential stability condition (29) and (30). In the former case, increasing N requires to increase M_R in order to keep $m_\nu = 0.2$ eV according to the equation (23). So for sufficiently high N the mass M_R runs over the condensation scale Λ . In the later case, the increase of N decreases the $\lambda_\nu(\mu)$ in an infrared region so that it eventually runs negative around the electroweak scale.

In Figs. 3 we plot M_h for various N from 1 to N^{\max} for three cases $\Lambda = 10^{16}, 10^{17}, 10^{18}$ GeV and we read out the intervals $(\mu_{t\nu}^{\min}, \mu_{t\nu}^{\max})$ of only possible values for $\mu_{t\nu}$ that correspond to $M_h = 125$ GeV. We show it in Tab. II together with the corresponding minimal and maximal masses for H^\pm , $(M_{H^\pm}^{\min}, M_{H^\pm}^{\max})$.

In Tab. II we show the maximum number N^{\max} as well. For the case $\Lambda = 10^{16}$ GeV ($\Lambda = 10^{17}$ GeV) the number

Λ [GeV]	$\mu_{t\nu}$ [GeV]	$\tan\beta$
10^{16}	44	1.183
10^{17}	54	1.215
10^{18}	62	1.245
10^{24}	86	1.401
10^{40}	101	1.624

TABLE I: Values of $\mu_{t\nu}$ and $\tan\beta$ depending on Λ while keeping $m_\nu = 0.2$ eV and $M_h = 125$ GeV for $N = 1$.

N^{\max} is actually not the maximal value allowed by either of conditions. It rather corresponds to maximizing the parameter $\mu_{t\nu}^{\max}$. Increasing N above N^{\max} causes the backward decrease of $\mu_{t\nu}$. The maximal number of right-handed neutrino triplets above which the non-decoupling condition is broken is $N = 158$ ($N \simeq 1500$). In the case $\Lambda = 10^{18}$ GeV, the maximum number $N^{\max} = 209$ is given by the Higgs potential stability.

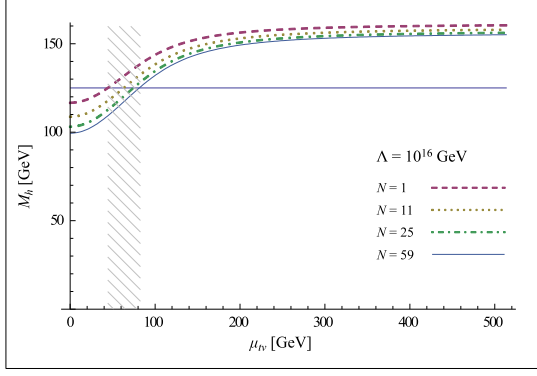
C. Light Higgs scalar coupling strengths

It is mandatory to ask how strongly does the candidate for the 125 GeV resonance, the lighter Higgs scalar h , couple to the fermions and gauge bosons. We study couplings relative to the Standard model case of h to W and Z bosons given by C_V , to top-quark given by C_t , and to charged Higgs bosons given by C_{H^\pm} , defined in (40). We plot their dependence on the number of right-handed neutrino triplets N in Fig. 4a) for three cases $\Lambda = 10^{16}, 10^{17}, 10^{18}$ GeV. The scaling coupling factor C_V approaches the Standard model value for larger N in the cases $\Lambda = 10^{17}$ GeV and $\Lambda = 10^{18}$ GeV. On the other hand, the coupling to top-quark given by C_t stays rather suppressed in comparison with the Standard model in the three cases.

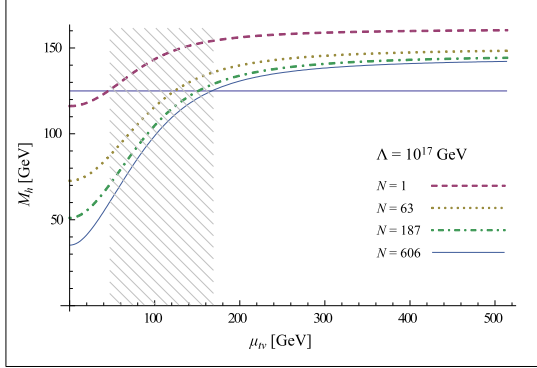
All three coupling parameters are relevant for the loop-induced decay of h to two photons. The dependence of the decay width $\Gamma(h \rightarrow \gamma\gamma)$ relative to the Standard model value $\Gamma(h \rightarrow \gamma\gamma)^{\text{SM}}$ on the number of right-handed neutrino triplets N is plotted in Fig. 4b). A slight enhancement occurs only for higher values of N for the cases $\Lambda = 10^{17}$ GeV and $\Lambda = 10^{18}$ GeV. The decay widths are calculated using the well known analytic expression to be found, e.g., in [33].

VII. DISCUSSION

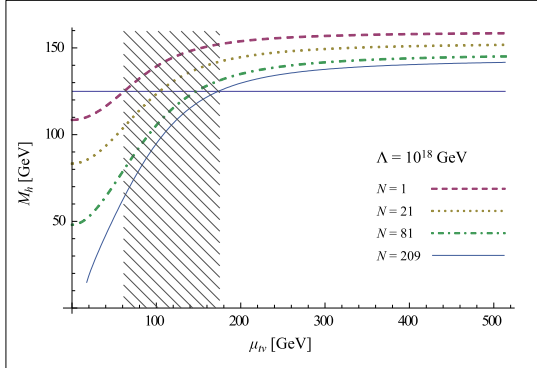
In this work we have chosen the two-Higgs doublet model as the next-to-simplest model accommodating the leading idea of our interest – identifying the top-quark and neutrino condensation as a sufficient source of the electroweak symmetry breaking – as it was proposed in [11].



a)



b)



c)

FIG. 3: The lightest scalar Higgs mass M_h for various numbers of right-handed neutrino triplets as a function of μ_{ν} for a) $\Lambda = 10^{16}$ GeV, b) $\Lambda = 10^{17}$ GeV, c) $\Lambda = 10^{18}$ GeV, and for fixed neutrino mass to be $m_\nu = 0.2$ eV. The solid horizontal line visualizes the 125 GeV value. The hatched area visualizes the interval of μ_{ν} values corresponding to $M_h = 125$ GeV shown in Tab. II as well.

Our analysis shows that it is possible to simultaneously reproduce correct values for top-quark mass m_t , the electroweak scale v , the 125 GeV boson mass, and the neutrino mass m_ν , below observational upper limit, without any unnatural parameter tension, despite rather limited manoeuvring space for participating parameters.

Λ [GeV]	N^{\max}	μ_{ν}^{\min} [GeV]	μ_{ν}^{\max} [GeV]	$M_{H^\pm}^{\min}$ [GeV]	$M_{H^\pm}^{\max}$ [GeV]
10^{16}	59	44	81	66	117
10^{17}	606	54	164	80	234
10^{18}	209	62	173	92	249

TABLE II: We show the maximum number of right-handed neutrino triplets N^{\max} for three cases $\Lambda = 10^{16}, 10^{17}, 10^{18}$ GeV. Next, we show the only intervals for μ_{ν} , $(\mu_{\nu}^{\min}, \mu_{\nu}^{\max})$, allowed by $M_h = 125$ GeV. Finally, we show the corresponding intervals for charged Higgs boson masses $(M_{H^\pm}^{\min}, M_{H^\pm}^{\max})$.

The number of right-handed neutrino types participating on the seesaw mechanism is not constrained phenomenologically by any upper limit [26, 27]. The model however exhibits an interesting feature providing us by such upper limit.

Next, we present two aspects of the model which are relevant for a present phenomenology: the mass spectrum of additional Higgs bosons H , A and H^\pm , and the coupling strengths of the 125 GeV Higgs boson to the top-quark and gauge bosons. We study their dependence on N and on the condensation scale Λ . Generally speaking, the higher values of N and Λ are preferred, because they lead to higher values of additional Higgs boson masses, and to the coupling strengths closer to the Standard model values. For example, for $\Lambda = 10^{18}$ and $N = 100$ we obtain the charged Higgs boson mass $M_{H^\pm} \doteq 223$ GeV and the coupling constant of h to W and Z at 93% level of the Standard model value.

The confrontation of the additional Higgs boson mass spectrum and the coupling properties of the lighter Higgs boson with the experimental constraints in the following two subsections however should be taken with a grain of salt for two reasons. First, the model analysed in this work is only a semi-realistic model: it ignores the mass generation of fermions other than top-quark and neutrinos and it is subject of simplification of the neutrino sector. Second, it is not possible to directly link our model to one of the standard types of two-Higgs-doublet models: we avoid a number of Higgs doublet Yukawa interactions to lighter fermions, and on the contrary we favour the coupling of one of the Higgs doublets to right-handed neutrinos.

A. Mass of charged Higgs boson

The values of M_{H^\pm} and $\tan \beta$ accessible in the model for the higher number of right-handed neutrino triplets N lie in the ranges

$$M_{H^\pm} \simeq (200 - 250) \text{ GeV}, \quad \tan \beta \simeq (1.2 - 1.25), \quad (55)$$

see Tab. I and Tab. II.

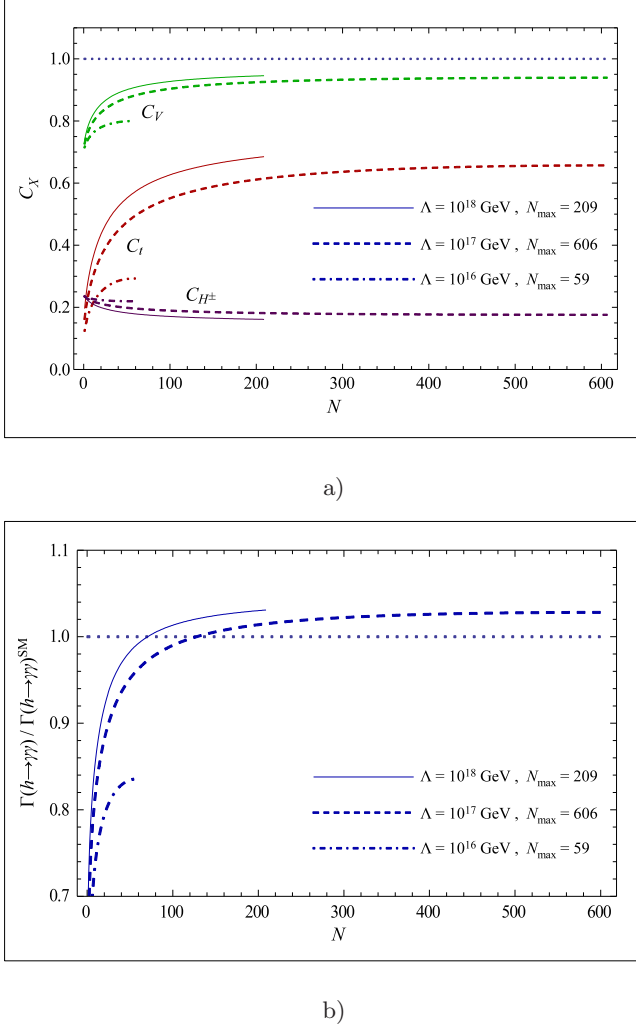


FIG. 4: For three cases $\Lambda = 10^{16}, 10^{17}, 10^{18}$ GeV we plot a) the N -dependence of relative coupling parameters C_t , C_V , and C_{H^\pm} defined in (40); b) the dependence of the decay width $\Gamma(h \rightarrow \gamma\gamma)$ relative to the Standard model value $\Gamma(h \rightarrow \gamma\gamma)^{\text{SM}}$ on N .

Direct searches for charged Higgs boson at LHC give a lower limit for its mass $M_{H^\pm} > 160$ GeV [34, 35]. The excluded mass interval corresponds to the below-threshold production $t \rightarrow H^+ b$. The analyses are made under assumption of 100% decay ratio via $H^+ \rightarrow \tau \nu$ and within one of the MSSM scenario. The lower limit for M_{H^\pm} translates in our model into the lower limit for N

$$N > 38 \quad \text{for} \quad \Lambda = 10^{17} \text{ GeV}, \quad (56)$$

$$N > 20 \quad \text{for} \quad \Lambda = 10^{18} \text{ GeV}, \quad (57)$$

while the case $\Lambda = 10^{16}$ GeV is excluded.

Indirect searches in B -physics are more stringent in setting the lower limits, but they are more model-dependent at the same time. For review see [36]. For example, the limit $M_{H^\pm} > 300 - 400$ GeV from the $B \rightarrow X_s \gamma$ decay is set for type-II two-Higgs-doublet model,

and the limit $M_{H^\pm} > 160$ GeV ($M_{H^\pm} > 500$ GeV)² from $Z \rightarrow b\bar{b}$ and from B - and K -meson mixing is set for type-I two-Higgs-doublet model. Even though it is rather speculative to apply those constraints to our model, they indicate that the model appears to be at the edge of viability.

B. Production and decays of lighter Higgs scalar

In order to successfully identify the lighter Higgs scalar h with the observed 125 GeV boson, it should exhibit coupling properties to other particles which leads to the observed phenomena.

Provided the higher number of right-handed neutrino triplets $N > 100$ and $\Lambda = 10^{18}$ GeV, the h scaling coupling factors (40) characteristic for the model are on the level

$$C_t \simeq (0.63 - 0.69), \quad C_V \simeq (0.93 - 0.95), \quad (58)$$

see Fig. 4a).

This result is to be compared with the ATLAS [37] and CMS [38] results for corresponding quantities.³ For example, the gluon fusion cross-section $gg \rightarrow h$ being induced by top-quark loop scales with a factor C_t^2 , or the partial decay width for $h \rightarrow WW$ scales with a factor C_V^2 . The best-fit values over all observed production-decay modes are

$$(C_t, C_V) \simeq (1.0, 1.2) \quad \text{ATLAS}, \quad (59)$$

$$(C_t, C_V) \simeq (0.5, 1.0) \quad \text{CMS}. \quad (60)$$

The point (58) lies within the ATLAS 95 % confidence level range and within the CMS less than 68 % confidence level range. The analyses were made under assumption that no non-Standard model particles contribute to the total decay width, what is reasonable assumption for our model, as they all are heavier than the lighter Higgs scalar h .

Out of the individual decay channels, we discuss $h \rightarrow \gamma\gamma$. Provided the higher number of right-handed neutrino triplets $N > 100$ and $\Lambda = 10^{18}$ GeV, the enhancement of the partial decay width $\Gamma(h \rightarrow \gamma\gamma)$ with respect to the Standard model can be achieved at the level of (1 - 3)%, see Fig. 4b),

$$\frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma(h \rightarrow \gamma\gamma)^{\text{SM}}} \simeq (1.01 - 1.03). \quad (61)$$

The signal strength μ for $h \rightarrow \gamma\gamma$ channel is measured with a 2σ excess with respect to the Standard model

² The limit in parenthesis follows from $B \rightarrow X_s \gamma$ but it is very sensitive to assumptions and to input parameters.

³ In the ATLAS and CMS analyses the overall fermion scaling factor C_F is used, instead of C_t , which scales only the top-quark Yukawa interaction in our model. In any case $C_t = C_F$.

[37, 38]

$$\mu(h \rightarrow \gamma\gamma) = \frac{\sigma_h}{\sigma_h^{\text{SM}}} \frac{\Gamma(h \rightarrow \gamma\gamma)/\Gamma_{\text{tot.}}}{\Gamma(h \rightarrow \gamma\gamma)^{\text{SM}}/\Gamma_{\text{tot.}}^{\text{SM}}} \sim (1.5 - 2.0). \quad (62)$$

The main part of the h production cross-section σ_h is given by gluon fusion cross-section $\sigma(gg \rightarrow h)$ which scales as $C_t^2 \simeq (0.39 - 0.47)$ in our model. This suppression can be however compensated by the suppression of the total decay width $\Gamma_{\text{tot.}}$, which scales as a linear combination of the C_t^2 and C_V^2 factors, both presenting a suppression. According to its parameter setting the model does not profit from the presence of charged Higgs boson in order to enhance significantly the $h \rightarrow \gamma\gamma$ signal strength with respect to the Standard model prediction.

VIII. CONCLUSIONS

We have analysed the top-quark and neutrino condensation scenario for the electroweak symmetry breaking within a reasonably simplified semi-realistic model with two composite Higgs doublets. We have demonstrated that it is possible to reproduce a mass spectrum of top-quark, neutrinos and observed bosons without any unnatural parameter tension. The lighter Higgs scalar is identified with 125 GeV particle. However some indirect constraints on the masses of yet unobserved additional Higgs bosons indicate that the model appears to be at the edge of viability. The coupling strengths of h differs from the Standard model values, but the experimen-

tal data have not a decisive power yet in this respect. If we keep only well established and model independent constraints then there remains some room for the top-quark and neutrino condensation scenario, provided that the condensation scale is $\Lambda \sim 10^{17-18}$ GeV and the number of right-handed neutrinos participating on the seesaw mechanism is $\mathcal{O}(100 - 1000)$.

There are two detail-independent predictions of the scenario. First, the h has rather big admixture of the neutrinos given by $\alpha \sim -0.8$. The mixing factor suppresses its Yukawa coupling with the top-quark and eventually with other charged fermions at the level of $\sim 60\%$ in comparison with the Standard Model. Second, the scenario provides an upper limit on the additional Higgs bosons which is rather low < 250 GeV. Through both predictions the scenario should be easily and definitely falsifiable by delivering more data from LHC in the near future.

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- [1] G. Aad et al. (ATLAS Collaboration), Phys.Lett.B (2012), 1207.7214.
 - [2] S. Chatrchyan et al. (CMS Collaboration), Phys.Lett.B (2012), 1207.7235.
 - [3] G. Degrandi, S. Di Vita, J. Elias-Miro, J. R. Espinosa, G. F. Giudice, et al., JHEP **1208**, 098 (2012), 1205.6497.
 - [4] A. Bednyakov, A. Pikelner, and V. Velizhanin (2013), 1303.4364.
 - [5] K. Chetyrkin and M. Zoller (2013), 1303.2890.
 - [6] J. Hošek (1985), preprint CERN-TH.4104/85.
 - [7] V. A. Miransky, M. Tanabashi, and K. Yamawaki, Phys. Lett. **B221**, 177 (1989).
 - [8] W. A. Bardeen, C. T. Hill, and M. Lindner, Phys. Rev. **D41**, 1647 (1990).
 - [9] S. P. Martin, Phys.Rev. **D44**, 2892 (1991).
 - [10] G. Cvetič and C. Kim, Mod.Phys.Lett. **A9**, 289 (1994).
 - [11] S. Antusch, J. Kersten, M. Lindner, and M. Ratz, Nucl.Phys. **B658**, 203 (2003), hep-ph/0211385.
 - [12] J. Hošek (1982), preprint JINR-E2-82-542, submitted to 21st Int. Conf. on High Energy Physics, Paris, France, Jul 26-31, 1982.
 - [13] K. Kimura and H. Munakata (1984).
 - [14] Y. Nagoshi, K. Nakanishi, and S. Tanaka, Prog. Theor. Phys. **85**, 131 (1991).
 - [15] G. Cvetič and C. Kim, Int.J.Mod.Phys. **A9**, 1495 (1994).
 - [16] V. Gribov, Phys.Lett. **B336**, 243 (1994), hep-ph/9407269.
 - [17] J. D. Bashford (2003), hep-ph/0310309.
 - [18] T. Brauner and J. Hosek (2004), hep-ph/0407339.
 - [19] P. Benes, T. Brauner, and A. Smetana, J.Phys. **G36**, 115004 (2009), 0806.2565.
 - [20] C. Wetterich, Phys.Rev. **D74**, 095009 (2006), hep-ph/0607051.
 - [21] J.-M. Schwindt and C. Wetterich, Phys.Rev. **D81**, 055005 (2010), 0812.4223.
 - [22] J. Hošek (2009), 0909.0629.
 - [23] J. Hošek (World Scientific, 2011), pp. 191–197, in proceedings of the Workshop in Honor of Toshihide Maskawa's 70th Birthday and 35th Anniversary of Dynamical Symmetry Breaking in SCGT, Nagoya University, Japan, 8-11 December 2009, Edts: Fukaya, H., Harada, M., Tanabashi, M., Yamawaki, K., World Scientific.
 - [24] P. Beneš, J. Hošek, and A. Smetana (2011), 1101.3456.
 - [25] A. Smetana (2011), 1104.1935.
 - [26] J. R. Ellis and O. Lebedev, Phys.Lett. **B653**, 411 (2007), 0707.3419.
 - [27] J. Heeck, Phys.Rev. **D86**, 093023 (2012), 1207.5521.
 - [28] M.-T. Eisele, Phys.Rev. **D77**, 043510 (2008), 0706.0200.
 - [29] B. Feldstein and W. Klemm, Phys.Rev. **D85**, 053007

- (2012), 1111.6690.
- [30] M. A. Luty, Phys.Rev. **D41**, 2893 (1990).
 - [31] C. T. Hill, C. N. Leung, and S. Rao, Nucl.Phys. **B262**, 517 (1985).
 - [32] G. Cvetič, Rev.Mod.Phys. **71**, 513 (1999), hep-ph/9702381.
 - [33] A. Djouadi, Phys.Rept. **459**, 1 (2008), hep-ph/0503173.
 - [34] S. Chatrchyan et al. (CMS Collaboration), JHEP **1207**, 143 (2012), 1205.5736.
 - [35] G. Aad et al. (ATLAS Collaboration), JHEP **1206**, 039 (2012), 1204.2760.
 - [36] G. Branco, P. Ferreira, L. Lavoura, M. Rebelo, M. Sher, et al., Phys.Rept. **516**, 1 (2012), 1106.0034.
 - [37] ATLAS Collaboration, 1204290 (2012).
 - [38] 1199142 (2012).